

Field theory of nucleon to higher-spin baryon transitions

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We discuss the nucleon to higher-spin N - and Δ -resonance transitions by pions and photons. The higher-spin baryons are described by Rarita-Schwinger fields and, as we argue, this imposes a stringent consistency requirement on the form of the couplings. Popular $\pi N\Delta$ and $\gamma N\Delta$ couplings are inconsistent from this point of view. We construct examples of consistent interactions with the same non-relativistic limit as the conventional ones.

11.10.Ef, 11.15.-q, 13.75.Gx, 14.20.Gk

Current experimental efforts promise to greatly advance the understanding of the strong and electroweak structure and the in-medium properties of the nucleon and its excited N^* -states, such as the spin-3/2 Δ -isobar [1–3]. In theoretical studies, an important role is played by relativistic Lagrangians describing on a hadronic level the interactions of the baryons with pions and photons. They are often used, at tree level or in a unitarized model, to extract the relevant coupling constants from the data, which can then be compared to results obtained in quark or other “QCD-inspired” microscopic models. To do this in a meaningful manner, one needs consistent interactions between nucleons, N^* -resonances, and pions and photons.

However, modern treatments of higher-spin ($s \geq 3/2$) baryon fields within the standard Rarita-Schwinger (RS) formalism are problematic. The difficulties are generic to any field-theoretic description of higher-spin particles and are related to the fact that (in a relativistic, local formulation in four space-time dimensions) a higher-spin field contains more components than is needed to represent the spin degrees of freedom (DOF) of the particle. The standard free-field formulations [4–6] are given by Lagrangians which, in addition to the Dirac- or Proca-type equations, yield constraint equations that reduce the number of independent components of the field to the correct value. The issue is how to introduce interactions: When these are not constructed consistently with the free theory, the constraints may be violated and consequently the unphysical extra DOF will become involved. The widely-used $\pi N\Delta$ and $\gamma N\Delta$ interactions given below in Eqs. (16) and (17) are examples of such inconsistent couplings. The pathologies of the $\pi N\Delta$ coupling have been especially thoroughly discussed [7–11]. It is the purpose of this Rapid Communication to present a remedy for these problems, construct explicit examples of alternative consistent interactions, and illustrate some consequences.

An elegant general way to distinguish consistent theories for high-spin fields is to use the correspondence between the local symmetries and the DOF content of the theory (see the theorem quoted below). The free massless theory can be con-

structed by demanding the action to be invariant under a number of gauge transformations [12–14], constraining the number of DOF to two. The mass term breaks these symmetries such that the number of DOF is raised to the appropriate $2s + 1$. Our basic premise is that a consistent interaction should not “activate” the spurious DOF, and therefore the full interacting theory must obey similar symmetry requirements as the corresponding free theory.

It is sometimes possible to formulate interactions which destroy the symmetries of the massless free theory, but in a way similar to the mass term. For instance, the “minimal” electromagnetic coupling of the RS field, or the conventional $\pi N\Delta$ coupling of Eq. (16) with the specific choice $z_\pi = 1/2$, are interactions of this type. It then appears that the constraints can be violated only for specific values of the interaction strength. However, such theories in general have non-positive-definite commutators – the Johnson-Sudarshan problem [15,8], as well as acausal propagations – the Velo-Zwanziger problem [16,9]. Also, correct derivations of Feynman rules in these theories indicate that Lorentz invariance is not obvious [17,11] despite the fact that one starts from a manifestly Lorentz-invariant Lagrangian. We shall not further discuss these problems in here, but merely assume that they should be absent in a consistent theory. We therefore adopt the viewpoint that consistent interactions must support the local symmetries of the free massless RS formulation, while the mass term breaks these symmetries in the correct manner.

In first instance, the interactions can be chosen to simply preserve the gauge symmetries of the free massless theory. In fact, the possibility to construct consistent higher-spin field theories with such gauge-invariant (GI) couplings was pointed out by Weinberg and Witten already some time ago [18], but apparently has never been exploited in hadronic physics. Here we shall confine ourselves to exploring this road, which is sufficient for formulating consistent $N \xrightarrow{\pi,\gamma} N^*$ interactions.

To present our arguments more systematically, let us first recall that within the RS formalism a field of spin $s = j + 1/2$ (j is an integer) is represented by a symmetric Lorentz tensor-spinor $\psi_{\mu_1 \dots \mu_j}^{(\alpha)}(x)$ of rank j ; the spinor index α will be omitted in what follows. Note that such a field has

$$C_j \equiv 4(j+1)(j+2)(j+3)/6 \quad (1)$$

independent components. The requirement that the field describes a massless particle, with only two helicities, leads to an essentially unique definition of the theory [6,12–14]. Namely, the action must be invariant under the gauge transformation

$$\delta \psi_{\mu_1 \dots \mu_j} = (1/j) [\partial_{\mu_1} \epsilon_{\mu_2 \dots \mu_j} + \dots + \partial_{\mu_j} \epsilon_{\mu_1 \dots \mu_{j-1}}] \quad , \quad (2)$$

where $\epsilon(x)$ is a symmetric tensor-spinor field of rank $j - 1$, subject to the traceless condition: $\gamma^{\mu_1} \epsilon_{\mu_1 \mu_2 \dots \mu_{j-1}} = 0$. Furthermore, the field itself must satisfy the gauge-invariant condition

$$g^{\mu_1 \mu_2} \gamma^{\mu_3} \psi_{\mu_1 \mu_2 \mu_3 \dots \mu_j} = 0. \quad (3)$$

To count the number of DOF we may use the Hamiltonian framework [19–22], where, given the Lagrangian density for ψ , one defines the conjugate momentum $\pi^{\mu_1 \dots \mu_j} = \partial \mathcal{L} / \partial \dot{\psi}_{\mu_1 \dots \mu_j}$ and determines all the constraints in the (π, ψ) phase space of the theory. Taking into account the condition Eq. (3), the field actually has

$$N_{\text{comp}} = C_j - C_{j-3} = 6j(j+1) + 4 \quad (4)$$

independent components, and so does its conjugate momentum. However, only two (for each of them) are needed to describe a massless particle. The rest is to be eliminated by means of the phase-space constraints. Let N_I and N_{II} denote the number of first- and second-class constraints, respectively. Each first-class (second-class) constraint eliminates two (one) DOF [21]. Thus, in a theory with only physical DOF N_I and N_{II} must satisfy

$$2N_{\text{comp}} - 2N_I - N_{II} = 2 + 2. \quad (5)$$

An explicit determination of the constraints in the most general case by, for instance, the usual Dirac-Bergmann procedure [19] is a formidable task. But their number can easily be assessed by using the following theorem, which establishes a precise correspondence between the local symmetries of the action and the first-class constraints:

Theorem. A Lagrangian theory invariant under a local transformation with n independent parameters has n primary first-class constraints [20]; and the total number of first-class constraints is $(d+1) \times n$, where d is the highest order of the time-derivative operator acting on the parameters of the transformation [23].

In our case, the Lagrangian is invariant under Eq. (2), hence $d = 1$ and $n = C_{j-1} - C_{j-2} = 2j(j+1)$, while

$$N_I = 4j(j+1) \quad (6)$$

by the second part of the theorem. Furthermore, since fermionic theories are of first order in space-time derivatives, the total number of primary constraints is equal to the number of field components, $N_I^{(1)} + N_{II}^{(1)} = N_{\text{comp}}$. Invoking the first part of the theorem we have $N_I^{(1)} = n$. From the fact that (in the massless case) there are only primary second-class constraints, i.e., $N_{II} = N_{II}^{(1)}$, we find

$$N_{II} = N_{\text{comp}} - n = 4j(j+1) + 4. \quad (7)$$

By using these values for N_I and N_{II} , one can check that Eq. (5) is indeed satisfied. We have thus proven the unitarity, or the so-called “no-ghost” theorem [6,12], of the massless higher-spin fermion formulation. A similar proof applies to the formulation for higher-spin bosons.

The mass term is usually introduced so as to break the gauge symmetry, turning all the first-class constraints into the second class. The resulting number of the second-class constraints \bar{N}_{II} must provide the physical DOF counting

$$2N_{\text{comp}} - \bar{N}_{II} = 2(2s+1). \quad (8)$$

One can in general write $\bar{N}_{II} = N_I + N_{II} + N'_{II}$, and find from Eqs. (6)–(8) that $N'_{II} = 4j^2$. This shows that the mass term should play a rather subtle role: besides turning the first-class constraints of the massless theory into the second class, some number N'_{II} of new second-class constraints must be generated.

The couplings consistent with the above free-theory construction, for both massless and massive cases, will apparently be only those which are invariant under transformation Eq. (2) or its modifications, the so-called “deformations,” with the same d and n . Realizations based on deformations of the free-theory symmetries appear to be unavoidable in the construction of “minimal” couplings of the RS fields to the photon or gravity. However, it is not necessary for constructing consistent $N \rightarrow N^*$ transition interactions. In this case, one can generally construct couplings invariant under Eq. (2).

To illustrate all this, let us specify the discussion to the spin-3/2 case, relevant to the important example of the Δ -isobar. The spin-3/2 field is described by the sixteen-component vector-spinor $\psi^\mu(x)$, with for the massless case the Lagrangian density

$$\mathcal{L} = \bar{\psi}^\lambda \mathcal{O}_{\lambda\varrho}(a) \frac{1}{2} \{ \sigma^{\varrho\mu}, i\partial \} \mathcal{O}_{\mu\nu}(a) \psi^\nu, \quad (9)$$

where

$$\mathcal{O}_{\mu\nu}(a) \equiv \exp(\frac{1}{4} a \gamma_\mu \gamma_\nu) = g_{\mu\nu} + \frac{1}{4} (e^a - 1) \gamma_\mu \gamma_\nu, \quad (10)$$

and the arbitrary constant a represents the freedom due to the point-transformation invariance [7,10]; $\sigma^{\varrho\mu} = [\gamma^\varrho, \gamma^\mu]/2$. The action of this theory is invariant under the gauge transformation

$$\delta\psi_\mu = \mathcal{O}_{\mu\nu}(-a) \partial^\nu \epsilon, \quad (11)$$

where $\epsilon(x)$ is a spinor field. By using the theorem we have $N_I = 8$ and $N_{II} = 12$, which is of course in agreement with an explicit evaluation of the constraints [11,24].

For $|a| < \infty$, the tensors $\mathcal{O}(a)$ form a group with $\mathcal{O}_{\mu\nu}(0) = g_{\mu\nu}$ as the unit element and product law $\mathcal{O}_\mu^{\varrho}(a_1) \mathcal{O}_{\varrho\nu}(a_2) = \mathcal{O}_{\mu\nu}(a_1 + a_2)$. Different choices among these finite values of a amount to the field redefinition $\psi'_\mu = \mathcal{O}_{\mu\nu}(a) \psi^\nu$. Since $\det \mathcal{O}(a) = e^a$ is a constant, any choice can be made without affecting the S matrix.

For the “forbidden” value $a = -\infty$, the Lagrangian becomes

$$\mathcal{L} = i\bar{\psi}^\lambda \mathcal{O}_{\lambda\mu}(-\infty) \gamma_\varrho \mathcal{O}^{\mu\nu}(-\infty) \partial^\varrho \psi_\nu, \quad (12)$$

where $\mathcal{O}_{\mu\nu}(-\infty) = g_{\mu\nu} - \frac{1}{4} \gamma_\mu \gamma_\nu$. This is the massless version of the theory recently considered by Haberkott [25]. Determining the constraints for this case we find $N_I = 4$ and

$N_{II} = 12$, hence more DOF than the RS theory. This can be understood by observing that the Lagrangian Eq. (12) is invariant under $\delta\psi_\mu = \gamma_\mu\epsilon$, *i.e.*, it has the same number of local transformations as the RS theory, but without the space-time derivative. Note that the massive case of this $a = -\infty$ theory [25] has the same $\delta\psi_\mu = \gamma_\mu\epsilon$ symmetry.

The massive RS theory is obtained by the replacement $\partial^\mu \rightarrow \partial^\mu + \frac{1}{4}iM\gamma^\mu$ in Eq. (9). The mass term breaks the gauge symmetry Eq. (11) in the correct way, raising the number of DOF to four as is appropriate for a massive spin-3/2 particle. The propagator of the theory is the well-known RS propagator; in terms of spin-projection operators $P^{(J)}$ [6,26] it reads¹

$$S_{\mu\nu}(p) = \frac{1}{\not{p} - M + i\epsilon} P_{\mu\nu}^{(3/2)} - \frac{2}{3M^2} (\not{p} + M) P_{22,\mu\nu}^{(1/2)} + \frac{1}{\sqrt{3}M} \left(P_{12,\mu\nu}^{(1/2)} + P_{21,\mu\nu}^{(1/2)} \right), \quad (13)$$

where

$$P_{\mu\nu}^{(3/2)} = g_{\mu\nu} - \frac{1}{3}\gamma_\mu\gamma_\nu - \frac{1}{3p^2}(\not{p}\gamma_\mu p_\nu + p_\mu\gamma_\nu\not{p}) \quad (14)$$

projects onto the pure spin-3/2 states, while

$$\begin{aligned} P_{22,\mu\nu}^{(1/2)} &= p_\mu p_\nu / p^2, \\ P_{12,\mu\nu}^{(1/2)} &= p^\rho p_\nu \sigma_{\rho\mu} / (\sqrt{3}p^2), \\ P_{21,\mu\nu}^{(1/2)} &= p_\mu p^\rho \sigma_{\rho\nu} / (\sqrt{3}p^2), \end{aligned} \quad (15)$$

are projection operators onto the spin-1/2 sector of the RS theory. The pole part of the RS propagator is proportional to $P^{(3/2)}$, while the nonpole part involves the spin-1/2 sector.

Consider next the interactions. In the literature, a popular choice for the $\pi N\Delta$ coupling is²

$$\mathcal{L}_{\pi N\Delta} = (f_{\pi N\Delta}/m_\pi) \bar{\psi}^\mu \Theta_{\mu\nu}(z_\pi) \Psi \partial^\nu \phi + \text{H.c.}, \quad (16)$$

while for the $\gamma N\Delta$ couplings one often takes

$$\begin{aligned} \mathcal{L}_{\gamma N\Delta}^{(1)} &= \frac{ieG_1}{2m} \bar{\psi}^\mu \Theta_{\mu\nu}(z_{\gamma,1}) \gamma_\nu \gamma_5 \Psi F^{\mu\nu} + \text{H.c.}, \\ \mathcal{L}_{\gamma N\Delta}^{(2)} &= -\frac{eG_2}{(2m)^2} \bar{\psi}^\mu \Theta_{\mu\nu}(z_{\gamma,2}) \gamma_5 \partial_\nu \Psi F^{\mu\nu} + \text{H.c.}, \\ \mathcal{L}_{\gamma N\Delta}^{(3)} &= -\frac{eG_3}{(2m)^2} \bar{\psi}^\mu \Theta_{\mu\nu}(z_{\gamma,3}) \gamma_5 \Psi \partial_\nu F^{\mu\nu} + \text{H.c.}; \end{aligned} \quad (17)$$

see, *e.g.*, Refs. [27–29]. Here, ψ^μ , Ψ , and ϕ denote the Δ -isobar vector-spinor, nucleon spinor, and pion pseudoscalar fields, with masses M , m , and m_π , respectively; $F^{\mu\nu}$ is the photon field tensor; $e \simeq \sqrt{4\pi/137}$ is the proton charge; $f_{\pi N\Delta}$ and G_i ($i = 1, 2, 3$) are dimensionless coupling constants. For real photons, only the G_1 and G_2 terms contribute.

The interactions in Eqs. (16) and (17) all contain the tensor $\Theta_{\mu\nu}(z) = g_{\mu\nu} - (z + \frac{1}{2})\gamma_\mu\gamma_\nu$; the constants z_π and $z_{\gamma,i}$ ($i = 1, 2, 3$) with arbitrary values are the so-called “off-shell parameters.”

These “conventional” $\pi N\Delta$ and $\gamma N\Delta$ interactions are inconsistent with the free spin-3/2 RS theory, for any value of the off-shell parameters, see, *e.g.*, Refs. [7–9,11]. They do not possess any local symmetries of the RS field, and as a consequence they violate the constraints and involve the unphysical lower-spin DOF. The latter contribute to the observables in terms of the “spin-1/2 backgrounds” [29].

In contrast, $N \xrightarrow{\pi,\gamma} \Delta$ couplings which are invariant under the gauge transformation Eq. (11) will be fully consistent in that sense. Such GI couplings can easily be constructed by using the manifestly invariant RS field tensor

$$G^{\mu\nu} = \partial^\mu \psi^\nu - \partial^\nu \psi^\mu \quad (18)$$

and its dual $\tilde{G}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma} G_{\rho\sigma}$. The corresponding vertices $\Gamma^\mu(p, p-k, k)$, where p and k are the momenta of the Δ and, *e.g.*, the pion, while μ is the Lorentz index associated with the Δ -field, will satisfy

$$p_\mu \Gamma^\mu(p, p-k, k) = 0. \quad (19)$$

From Eqs. (13) and (15) one can then immediately see that all nonvanishing Δ -exchange amplitudes (see, *e.g.*, Fig. 1),

$$\Gamma^\mu(p, p-k', k') S_{\mu\nu}(p) \Gamma^\nu(p, p-k, k),$$

are proportional to the spin-3/2 projection operator, and thus the unphysical spin-1/2 sector decouples. (Let us remark here that it was correctly anticipated in Refs. [30,31] that Δ -exchange amplitudes should be proportional to the spin-3/2 projection operator. However, the *ad-hoc* prescriptions used in these works cannot be derived from a local Lagrangian and face other problems [10]. They are also at variance with the standard result [32] that the amplitude for exchange of a spin- s particle must behave as $p^{2s}/(p^2 - m^2)$. These criticisms do not apply to the present approach, where a typical Δ -exchange amplitude is given by Eq. (22) below.)

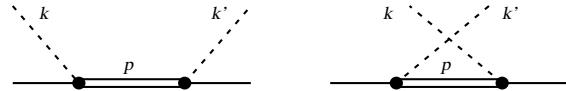


FIG. 1. Typical Δ -exchange graphs. Here the dashed, solid, and double lines denote the pion, the nucleon, and the Δ propagations, respectively.

In principle, there are many GI couplings one can construct. We will focus here only on the ones that become equivalent to the conventional couplings at the Δ -pole and hence have the same nonrelativistic limit. By using this “pole equivalence” we can establish the connection with coupling constants used in the vast number of previous studies, including the nonrelativistic [33] and heavy-baryon [34,35] formalisms.

¹In what follows, we choose $a = 0$ without loss of generality.

²For brevity, isospin space is omitted throughout.

For the $\pi N \Delta$ interaction we take

$$\mathcal{L}_{\pi N \Delta} = f \bar{\Psi} \gamma_5 \gamma_\mu \tilde{G}^{\mu\nu} \partial_\nu \phi + \text{H.c.} \quad (20)$$

The pole equivalence implies that the vertex Γ^μ obtained from this Lagrangian, when contracted with the free RS vector-spinor $u_\mu(\hat{p})$, where \hat{p} is the on-shell momentum of the Δ , $\hat{p}^2 = M^2$, becomes equivalent to the conventional $\pi N \Delta$ vertex, Γ_{conv}^μ , found from Eq. (16), *i.e.*,

$$\Gamma^\mu(\hat{p}, \hat{p} - k, k) u_\mu(\hat{p}) = \Gamma_{\text{conv}}^\mu(\hat{p}, \hat{p} - k, k) u_\mu(\hat{p}). \quad (21)$$

This condition requires us to identify the coupling constant in Eq. (20) as $f = f_{\pi N \Delta} / (m_\pi M)$, in terms of the coupling $f_{\pi N \Delta}$ of Eq. (16). However, we emphasize that despite the imposed pole equivalence, the two couplings will give different results for Δ -exchange amplitudes of, *e.g.*, Fig. 1, even at the pole $p^2 = M^2$. This is because with the conventional coupling one still encounters the background due to the negative-energy state contribution of the spin-1/2 sector. In contrast, using the GI interaction Eq. (20), we obtain the amplitude

$$\begin{aligned} \Gamma^\mu(p, p - k', k') S_{\mu\nu}(p) \Gamma^\nu(p, p - k, k) \\ = \frac{(f_{\pi N \Delta} / m_\pi)^2}{\not{p} - M} \frac{p^2}{M^2} P_{\mu\nu}^{(3/2)} k'^\mu k^\nu, \end{aligned} \quad (22)$$

for any Δ -momentum p . Some realistic calculations of πN scattering lengths and phase shifts using this amplitude have recently been reported [36,37]. These studies indicate large qualitative differences with the conventional approach, while in both approaches agreement with experiment can be achieved due to the interplay of other reaction mechanisms.

Considering the photon couplings, the GI $\gamma N \Delta$ interactions that are lowest in number of derivatives read

$$\begin{aligned} \mathcal{L}_{\gamma N \Delta} = e \bar{\Psi} \left(g_1 \tilde{G}_{\mu\nu} + g_2 \gamma_5 G_{\mu\nu} + g_3 \gamma_\mu \gamma^\rho \tilde{G}_{\rho\nu} \right. \\ \left. + g_4 \gamma_5 \gamma_\mu \gamma^\rho G_{\rho\nu} \right) F^{\mu\nu} + \text{H.c.} \end{aligned} \quad (23)$$

The first term contributes purely to the magnetic-dipole transition in the Sachs-type decomposition of the $\gamma N \Delta$ vertex [38]. The second term is up to a total derivative equal to the sum of the conventional G_2 and G_3 couplings of Eq. (17), provided that $G_2 = G_3 = (2m)^2 g_2$ and $z_{\gamma,2} = z_{\gamma,3} = -1/2$. Therefore, for real photons the g_2 coupling and the G_2 coupling with $z_{\gamma,2} = -1/2$ are fully equivalent. The g_3 and g_4 terms are new. However, at the Δ -pole the g_3 and G_1 couplings become equivalent, provided that $g_3 = G_1 / (2mM)$. The same applies to the g_4 term. Thus, the contribution of the GI couplings to the magnetic-dipole G_M and the electric-quadrupole G_E transition form factors at the Δ -pole reads,³ in the conventions of Ref. [38],

$$\begin{aligned} 3G_M &= 2m(M + m)g_1 - m(M - m)g_2 \\ &\quad + m(3M + m)(g_3 + g_4), \\ 3G_E &= m(M - m)(g_3 + g_4 - g_2). \end{aligned} \quad (24)$$

³A complete treatment of the Coulomb (longitudinal) quadrupole requires higher derivatives than in Eq. (23).

Since the considered $\gamma N \Delta$ couplings are invariant under both the electromagnetic and the RS gauge transformation, the corresponding vertex $\Gamma^{\mu\rho}(p, q)$, where q is the photon momentum, obeys the transversality condition with respect to both indices, *i.e.*,

$$p_\mu \Gamma^{\mu\rho}(p, q) = q_\rho \Gamma^{\mu\rho}(p, q) = 0. \quad (25)$$

Note that we have excluded couplings that contain $\sigma_{\mu\nu} G^{\mu\nu}$. Such couplings project onto the purely spin-1/2 contribution ($\gamma_\mu \gamma_\nu$ can be written in terms of the spin-1/2 projection operators only), which on the other hand decouple because of the gauge symmetry. Hence, these couplings lead to vanishing amplitudes, see Ref. [11] for an example.

An interesting extension is to study interactions that are not exactly gauge invariant, but the variation of which is proportional to some free-field equations. The invariance of the full action can then be provided by a variation of a corresponding free action, as illustrated below. In such theories, the decoupling of the unphysical DOF happens only when the particles, the free-field equations of which become involved, are on their mass shell. An analogous situation arises, for instance, in QED where the electromagnetic currents are conserved only when the external lepton legs are on-shell. This can be physically acceptable since the spurious DOF, even though present off-shell, do not contribute to observables. Moreover, such interactions may have the advantage of being lower in number of derivatives than the explicitly gauge-invariant ones.

However, for the $N \rightarrow \Delta$ case our attempts to find an interaction of this type led only to (locally) supersymmetric realizations. Consider, *e.g.*,

$$\mathcal{L} = g \bar{\Psi} \gamma_\mu (i \not{\partial} \phi + m \phi) \psi^\mu, \quad (26)$$

where for now all the fields are Hermitian. Under the variation $\delta\psi_\mu = \partial_\mu \epsilon$, we have, up to a total derivative,

$$\begin{aligned} \delta\mathcal{L} &= -g [(i \partial_\mu \bar{\Psi} \gamma^\mu + m \bar{\Psi}) (\not{\partial} \phi - i m \phi) \\ &\quad + i \bar{\Psi} (\partial^2 \phi + m^2 \phi)] \epsilon, \end{aligned} \quad (27)$$

which is indeed proportional to the free-field equation, if the pseudoscalar and the spinor field have the same mass equal to m . This variation is cancelled by the variation of the free Lagrangian, $\mathcal{L}_0 = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m^2 \phi^2 + \frac{1}{2} \bar{\Psi} (i \not{\partial} - m) \Psi$, under the local transformation

$$\delta\phi = i g \bar{\Psi} \epsilon, \quad \delta\Psi = g [\not{\partial} \phi - i m \phi] \epsilon,$$

which obviously is a supersymmetric transformation. It can still be suspected that this model is not fully consistent: a nontrivial supersymmetry necessitates the balance between fermionic and bosonic DOF, which may thus require inclusion of more boson fields as well as more interaction terms, see, *e.g.*, [26]. However, we shall not pursue this line here. As far as $N \rightarrow \Delta$ couplings are concerned, a development of such supersymmetric models seems neither necessary nor promising at present.

In conclusion, we have shown how the requirement of gauge invariance allows one to incorporate both manifest covariance and consistent DOF counting in a local high-spin

field formulation. We have therefore proposed to use the gauge-invariant interactions for describing the various meson- and photon-induced $N \rightarrow N^*$ transitions. This approach has been illustrated by the example of a spin-3/2 N^* -resonance, the Δ -isobar. While the conventional interactions used to describe the nucleon to Δ -isobar transitions by pions and photons are well-known to be inconsistent, we have constructed explicit examples of novel consistent interactions that become equivalent to the conventional ones at the Δ -pole and in the nonrelativistic limit. We emphasize that, even though most of our discussion has been focused on the interactions *explicitly* invariant under the free-field transformation Eq. (2), one can possibly construct consistent interactions based on its deformations with the same parameters. Moreover, interactions which are invariant up to some free-field equations can also be consistent, but in our attempts to construct such a $\pi N \Delta$ coupling we were led to a locally supersymmetric realization which has an obscure phenomenological implementation. On the other hand, our results concerning the explicitly gauge-invariant interactions are certainly relevant to theoretical studies related to the ongoing experimental programs on meson-, photo-, and electroproduction of N^* -resonances.

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